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A New Approach to the Phenomenology of Cosmological Supersymmetry Breaking

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Abstract: I introduce a new low energy effective description of Cosmological SUSY breaking. It requires the existence of a strongly interacting gauge theory at a scale of order 10^3 GeV, some of whose fields carry standard model quantum numbers. SUSY breaking is communicated to the standard model through gaugino masses and the Higgs multiplets. The model also provides a possible new solution of the strong CP problem.

Keywords: Cosmological SUSY Breaking.

Contents

1. Introduction

The hypothesis of Cosmological SUSY Breaking[1] (CSB) correlates the gravitino mass, $m_{3/2}$ with the cosmological constant, according to the formula

$$F_G \sim m_{3/2} M_P \sim \Lambda^{1/4} M_P. \tag{1.1}$$

Λ is viewed as a discrete, tunable parameter (perhaps determined in the real world by galactothropic considerations), and the *limiting model* with vanishing Λ is assumed to preserve exact $N = 1,\ d = 4$ super-Poincare invariance and a discrete R-symmetry. The scaling law (1.1) for the corrections to this limit, was originally postulated on phenomenological grounds. More recently, I provided a hand waving derivation of this result[2]. The origin of the large SUSY breaking effects is interaction with the horizon states of a stable dS space.

The low energy effective theory resulting from such a picture is highly constrained. First we must find an isolated super-Poincare invariant solution of string/M - theory.

We will call this the *limiting model*. Then, small explicit R-symmetry breaking perturbations of the low energy Lagrangian, must lead to spontaneous breaking of SUSY.

This can only occur if the limiting model at $\Lambda = 0$ has a massless fermion, ready to play the role of the Goldstino. In a previous paper[3] I gave arguments that the Goldstino must be a member of a vector superfield. Those arguments were incomplete, and the model of this paper is a counterexample. The models discussed in [3] suffered from a number of phenomenological problems. They also invoked a field independent Fayet-Iliopoulos term for a $U(1)$ gauge symmetry. Witten has argued[4] that such a term is inconsistent with the Dirac quantization condition.

In the present model, the Goldstino will be a chiral superfield G, which transforms as a singlet under all low energy gauge groups, and has vanishing R charge. Terms in the superpotential, depending only on G, must vanish in the limiting model. The G dependence of the superpotential for small Λ is generated by the mechanism described in[2]. It has the form:

$$\Lambda^{1/4} M_P^2 \; w(G/M_P). \tag{1.2}$$

We will arrange the rest of the dynamics of the model so that the VEV of $|G|$ is $\ll M_P$, and consequently $F_G \sim \Lambda^{1/4} M_P$.

2. Gaugino masses and a new low energy gauge group

The rest of our proposal for the dynamics of the G field is motivated by the phenomenological requirement of large gaugino masses. We will introduce a new low energy gauge group $\mathcal{G}$ with a variety of chiral multiplets, some of which also transform under the standard model gauge group. In order to preserve coupling unification, it is probably best to assign the standard model quantum numbers of complete $SU(5)$ multiplets to chiral supermultiplets transforming under $\mathcal{G}$.

This new gauge theory must satisfy the following criteria (we will discuss the problem of constructing explicit examples in a later section).

- It becomes strongly coupled at a scale M_1 but does not break either SUSY or R-symmetry.
- There is a gauge invariant, renormalizable Yukawa coupling $g_{\mathcal{G}} \int d^2\theta G F_1 F_2$, where the F_i transform under $\mathcal{G}$.
- Apart from a term $F_1 F_2$, which can be absorbed by a linear shift of G, the model has no allowed relevant operators. We choose the origin of G to eliminate the relevant operator.

- There is an R allowed coupling $g_\mu \int d^2\theta G H_u H_d$. The conventional μ term is forbidden by a symmetry $\mathcal{F}$, which we will introduce below.

When we integrate out scales above M_1, we obtain an effective action for G. The assumptions of unbroken R - symmetry implies that the part of the action which depends only on G consists of a Kahler potential of the form

$$K = GG^* k(\frac{G}{M_1}, \frac{G^*}{M_1}, \frac{G}{M_U}, \frac{G^*}{M_U}), \tag{2.1}$$

where M_U is the unification scale. Here we make the assumption that the coupling of G to any part of the spectrum at scales intermediate between M_1 and M_U is suppressed by at least an inverse power of M_U. It might even be reasonable to replace M_U with M_P in this formula. I believe that the answer to this question depends on details of the limiting model at scales of order M_U.

When combined with the superpotential (1.2), this Kahler potential gives a potential for $|G|$, which varies on the scale M_1. We will assume that it has a minimum at $< |G| >\sim M_1$. SUSY is broken at this minimum. The value of the F term of G is given approximately by

$$F_G = \sqrt{K^{GG^*}(< |G| > /M_1, < |G| >^* /M_1)} w'(0) \Lambda^{1/4} M_P \sim 10^7 (\mathrm{GeV})^2. \tag{2.2}$$

This gives rise to a gravitino mass of order 10^{-3} eV. The formula for the mass scales with the power of Λ predicted in [2]. The value of $w(0)$, which is a number of order 1, must be fine tuned to an accuracy $\frac{\Lambda^{1/2}}{M_P^2}$ in order to produce the correct value, Λ, for the value of the effective potential at its minimum. Λ is a fundamental input parameter in CSB, rather than a calculable low energy effective parameter, so this fine tuning is philosophically unexceptional. If one wishes, one can determine the correct value of this parameter in the real world, by applying the *galactothropic principle* of Weinberg[5], rather than simply fitting more recent cosmological data.

To get an estimate for what we want M_1 to be, we calculate the gaugino masses[1]

[1] This formula would follow from couplings of the form $(G/M_1)^a W^2$ or a Wess-Zumino coupling $\ln(G/M_1) W^2$, if $< |G| >\sim M_1$. The latter form would be natural if the $\mathcal{G}$ gauge theory has an accidental anomalous $U(1)$ symmetry, U_A, (with anomalies coming from the standard model), under which G transforms. For appropriate values of the anomaly coefficients, the WZ form preserves a discrete subgroup $\mathcal{F}$ of this $U(1)$. This is the symmetry we need to ensure the naturalness of the size of the μ term. On the other hand, if the $\mathcal{G}$ theory breaks U_A to $\mathcal{F}$ by either classical superpotential terms or a quantum anomaly, we could have couplings of the form $(G/M_1)^a W^2$, where G^a is the lowest dimension holomorphic $\mathcal{F}$ invariant, we can construct from G. If $< |G| >\sim M_1$ this gives a similar formula for gaugino masses. The WZ form can give us a QCD axion, a possibility we discuss below.

$$m_{1/2}^{(i)} \sim \frac{\alpha_i}{\pi}\epsilon_i \frac{F_G}{M_1}. \tag{2.3}$$

The running couplings α_i in this formula, are to be evaluated at the gaugino mass scale. ϵ_i are determined by calculations in the strongly coupled gauge theory at scale M_1. If $\epsilon_2 \sim 1$, this gives a wino mass of order 100 GeV if $M_1 \sim 1$ TeV. Thus, a reasonable value for gaugino masses requires a near coincidence between the dynamical scale M_1 and the scale of CSB, $\sqrt{M_P \Lambda^{1/4}}$.

The origin of the G field is determined by its coupling to the gauge theory $\mathcal{G}$, so that we are not allowed to simply absorb the conventional μ term into the VEV of G. However, if we introduce a discrete symmetry $\mathcal{F}$ under which G transforms by a phase, and assume that the coupling $\int d^2\theta \, G H_u H_d$ is $\mathcal{F}$ invariant, then the conventional μ term will be forbidden. Like the discrete R symmetry which guarantees Poincare invariance in the limiting model, $\mathcal{F}$ will be explicitly broken by interactions with the horizon. This breaking is sufficiently small to ignore. The dominant breaking of $\mathcal{F}$ will come from the VEV of G. $\mathcal{F}$ is required to be a symmetry of the $\mathcal{G}$ gauge theory.

3. Baryon number, lepton number, and flavor

A central element in CSB is the discrete R symmetry which guarantees Poincare invariance in the the limiting model. This can be put to other uses. Here we will show that it can eliminate all unwanted dimension 4 and 5 baryon and lepton number violating operators in the supersymmetric standard model. The dimensionless coefficients of these operators will thus be suppressed by at least $\frac{\Lambda^{1/8}}{\sqrt{M_P}} \sim 10^{-15.5}$. This is sufficient to account for experimental bounds on baryon and lepton number violating processes. The interaction $\int d^2\theta \; H_u^2 L^2$, will need a separate discussion. It should not be forbidden by R.

We will choose the R charge of SSM fields to be independent of quark and lepton flavor, and denote it by the name of the corresponding field. All R charges are to be understood modulo N, where Z_N is the R symmetry group. Flavor dependent R charges would require many important Yukawa couplings to vanish, and the corrections to the R symmetric limit are too small to account for the non-zero values of these couplings.

The condition that the standard Yukawa couplings are allowed by R symmetry is

$$L + H_d + \bar{E} = Q + H_d + \bar{D} = Q + H_u + \bar{U} = H_u + H_d = 2. \tag{3.1}$$

Note that, although these conditions allow a term $\int d^2\theta \; H_u H_d$, it will be forbidden by $\mathcal{F}$. We will also impose $2L + 2H_u = 2$ to allow the dimension 5 F term which can

generate neutrino masses. The renormalizable dynamics of the $\mathcal{G}$ gauge theory, must have an accidental symmetry which forbids the generation of this term with coefficient $\frac{1}{M_1}$. The combination of the accidental symmetry and the group $\mathcal{F}$ introduced below, should also forbid other dimension 5 lepton number violating operators (both D and F terms) with a coefficient of this scale. Neutrino masses can then be generated by dynamics at the scale M_U[2]

Dimension 4 baryon and lepton number violating operators in the superpotential will be forbidden in the limiting model by the inequalities

$$2L + \bar{E} \neq 2 \tag{3.2}$$

$$2\bar{D} + \bar{U} \neq 2, \tag{3.3}$$

$$L + Q + \bar{E} \neq 2. \tag{3.4}$$

Absence of dimension 5 baryon number violating operators requires

$$3Q + L \neq 2 \tag{3.5}$$

$$3Q + H_d \neq 2 \tag{3.6}$$

$$\bar{E} + 2\bar{U} + \bar{D} \neq 2, \tag{3.7}$$

The condition that there be no baryon number violating dimension 5 D-terms is that none of $Q + \bar{U} - L$; or $U + E - D$, vanishes.

These equations can be simplified by solving the equalities for $\bar{E}$, $\bar{U}$, $\bar{D}$, and H_u in terms of Q, L, and H_d. The conditions then become

$$2 \neq 3Q + H_d \neq 3 \tag{3.8}$$

$$2Q + H_d \neq 5, \tag{3.9}$$

. Recall that all of these conditions are to be understood modulo $N \geq 3$, where Z_N is the discrete R symmetry.

A possible solution of all of these constraints is $H_d = 1$, $Q = 0$, $L = 0$, $H_u = 1$, $\bar{E} = 1$, $\bar{D} = 1$, $\bar{U} = 1$. The discrete group can be Z_N with $N \geq 3$.

Thus, if the discrete R symmetry group and its representations in the SSM are chosen appropriately we can understand both the absence of unacceptable baryon and lepton number violating operators, and the presence of neutrino masses. The observed

[2]It is well known that this gives neutrino masses an order of magnitude too small to explain experiment. We have no found no neat solution to this problem in the present context. Later we speculate that this factor may be related to the small numbers which appear in the quark mass matrix, and might be explained by the Froggatt-Nielsen mechanism.

size of neutrino masses puts a constraint on the new physics at the scale M_1 . There must be an accidental symmetry of the combined $\mathcal{G}\times SU(3)\times SU(2)\times U(1)$ gauge theory, which forbids the lepton number violating dimension five operators that could lead to neutrino masses of order $\frac{<H_U>^2}{M_1}$. The exact symmetries R and $\mathcal{F}$ must permit this dimension 5 operator. They are broken by effects on the horizon, but these mechanisms would induce this operator with a coefficient much too small to account for neutrino masses. We expect neutrino masses to be determined by physics near the GUT scale, which generates this operator with coefficient $\sim \frac{10}{M_U}$.

3.1 Flavor

The simplest solution of flavor problems in this model is to assume that the origin of flavor breaking is in physics near the GUT scale. The low energy theory at scales $\sim M_1$ is a $\mathcal{G} \otimes SU(3,2,1)$ gauge theory. It has a large flavor symmetry acting on quarks and leptons, which is broken only by the standard Yukawa couplings of H_u and H_d. If all other physical excitations have masses near the unification scale, the GIM mechanism is operative and flavor changing processes occur at acceptable levels.

In a more ambitious model, the Goldstino field G might also allow us to implement the Froggatt-Nielsen (FN) [6] mechanism for explaining the flavor structure of quark and lepton masses and mixing angles. The basic idea is very simple. We postulate that the discrete symmetry group $\mathcal{F}$, which commutes with the supercharges, and under which G transforms by a phase, also acts on quarks and leptons in a flavor dependent way[3]. This symmetry is of course broken for finite values of the cosmological constant, but this explicit breaking is much smaller than the spontaneous breaking due to the VEV of G. In this version of the model, the $\mathcal{F}$ charges are family quantum numbers which distinguish quark and lepton flavors.

In order to account for quark and lepton mass matrices of appropriate size, we need, at the level of the effective field theory containing G and the standard model, nonrenormalizable couplings between G and standard model chiral multiplets, which are scaled by M_1 rather than the Planck mass or unification scale. The non-renormalization theorem for superpotentials, and our assumption of unbroken R - symmetry of the $\mathcal{G}$ dynamics, will require us to have Yukawa couplings between standard model chiral fields, and the fields which are charged under $\mathcal{G}$. These terms could have the form $\int d^2\theta \; T\bar{T}^A\bar{5}_A$, where T is in some $\mathcal{G}$ representation, and a standard model singlet. $\bar{T}^A$ transforms in the conjugate representation of $\mathcal{G}$ and the [5] of $SU(5)$, while $\bar{5}_A$ are the usual standard model fields which fit into the [$\bar{5}$] of $SU(5)$. Of course, at this

[3]It must act on the Higgs fields in order to explain the value of the μ term. It must then act on quarks and leptons if the standard Yukawa couplings are to preserve $\mathcal{F}$. However, we have a choice of whether this action is flavor blind or not.

low energy level, $SU(5)$ is broken, and there is no reason for these Yukawa couplings to satisfy $SU(5)$ relations (though we have only written the $SU(5)$ invariant coupling explicitly).

When we integrate out physics at the scale M_1, these couplings will give rise to irrelevant couplings in the effective theory which describes G and the MSSM. In particular, there will be terms in the superpotential of the form $\int d^2\theta \ \lambda_u^{ij}(G/M_1)H_u Q_i \bar{U}_j$ with similar terms for down quarks and leptons. We assume that the $\mathcal{G}$ dynamics does not break either $\mathcal{F}$ or the R-symmetry, so[4] these matrix valued functions of G/M_1 must respect these symmetries. In particular, a given power of G can only appear if its $\mathcal{F}$ quantum numbers are neutralized by those of the quarks or leptons. Assuming that $< G > /M_1 \sim .2$[5], we get a Yukawa coupling matrix whose entries are powers of this small parameter. This is the Froggatt-Nielsen mechanism. There are a large number of papers on the Froggat-Nielsen mechanism[7], and it is well known that it is possible to construct models of this type which give correct predictions for the quark and lepton masses and the CKM angles.

However, Y. Nir[9] has informed me that it is very difficult to make phenomenologically consistent models of this type at low scales. Thus, this superficially attractive possibility will probably lead to phenomenological problems. It is likely then that the theory of flavor, like that of neutrino masses, is associated with scales of order the unification scale. Note by the way that the Froggatt-Nielsen mechanism for explaining flavor hierarchies, requires two closely spaced energy scales in order to account for the small parameter whose powers govern the texture of the quark and lepton mass matrices. It is interesting to speculate that it is a power of the same small parameter which accounts for the otherwise mysterious order of magnitude discrepancy between the neutrino see-saw scale and the unification scale. That is, the neutrino see-saw scale might be $\theta^b M_U$, with $b = 1$ or 2 and θ the Cabibbo angle.

4. $SU(2) \otimes U(1)$ breaking

There is no reason for either of the Yukawa couplings $g_H \int d^2\theta G H_u H_d$ or $g_{\mathcal{G}} \int d^2\theta G F_1 F_2$ to be particularly small. If we integrate out degrees of freedom above the scale M_1,

[4]apart from tiny terms coming from interaction with the horizon,

[5]If we have the WZ coupling of G to the standard model gauge fields, which would follow from an unbroken accidental U_A symmetry of $\mathcal{G}$ physics, this assumption also increases our estimate of gaugino masses by a factor of 5, thus raising the scale M_1. If, on the other hand, we have couplings of the form $G^a W^2$, then the small VEV of G appears in the numerator of our estimate for the gaugino masses. It is no longer possible to have a viable effective field theory. Thus, we can only have a FN mechanism based on G, if the phase of G is a QCD axion. We will see below that this is probably ruled out experimentally.

the first of these couplings produces Kahler potential terms of the form $\frac{1}{M_1} G G^* H_u H_u^*$ and $\frac{1}{M_1} G G^* H_d H_d^*$. These terms are of order $\frac{g_H^4}{4\pi^2}$. There are then strong $\mathcal{G}$ corrections to this suppressed by a further factor of $\frac{g_{\mathcal{G}}^2}{4\pi^2}$. If the couplings are not small these terms produce significant contributions to the quadratic term in the Higgs potential, when we insert the F term of G. The sign of the quadratic terms may depend on the details of strongly interacting physics at scale M_1, if $g_{\mathcal{G}}$ is not small. If the sign is negative and dominates other contributions to the potential, we find the required breakdown of the weak interaction gauge symmetry. Of course, we can also get a contribution of the same type, from the top quark Yukawa coupling, as in gauge mediation.

Notice that F_G produces tree level SUSY breaking in the Higgs supermultiplets. Higgs loops will then give the dominant contribution to the splitting between the top quark and top squark (and perhaps the bottom quark/squark splitting as well if m_t/m_b is attributed to large $\tan\beta$.). Other squarks and sleptons will get their masses predominantly through gauge loops, as in gauge mediated models.

Thus, we can expect the spectrum of gauginos, sleptons and light squarks to resemble that of gauge mediated models, while squark partners of the heavy quarks, and particles in the Higgs multiplets will have masses which depend on the new Yukawa couplings g_H and $g_{\mathcal{G}}$, and, if the latter is strong enough, also on the details of strong $\mathcal{G}$ dynamics. More work is necessary to determine whether there are any potential problems with existing measurements, and to sharpen the predictions of the model for physics accessible to the LHC.

4.1 Dark matter

The gravitino is the lightest fermion in this model, and will be stable. It is relatively strongly coupled to the rest of the system, through its Goldstino component G. Thus, there will not be a WIMP LSP candidate for dark matter. The most likely dark matter candidate is a cosmologically stable $\mathcal{G}$ hadron. The $\mathcal{G}$ theory is required to have a variety of exact and accidental symmetries to account for the scale of the μ parameter, neutrino masses, *etc.* It would not be surprising to find that these implied a quasi-stable particle whose mass and annihilation cross section were related to the scale M_1. More detailed analysis will be required to determine if such a particle is a viable dark matter candidate, but it is in the right ballpark[6]

4.2 A QCD axion and an alternative solution to the strong CP problem

The phase of G is an angular variable, which might couple to QCD like a Peccei-Quinn-

[6]Note that the mass is a bit heavier than conventional WIMPs, but the annihilation cross section is likely to be larger because it does not contain weak dimensionless couplings.

Weinberg-Wilczek axion, if U_A symmetry is unbroken by $\mathcal{G}$ dynamics. However, the axion decay constant is in a range which is almost certainly ruled out by experiment. In our low energy model, at scale M_1 we can postulate an accidental axial symmetry which acts on G, and guarantees that its Kahler potential has the form $K(G\bar{G}/M_1^2)$, and that the couplings to standard model gauge bosons have the WZ form. In this approximation, G is a QCD axion with decay constant $\sim M_1$. Dynamics at and above the unification scale has no reason to preserve this symmetry. It need only preserve the discrete subgroup $\mathcal{F}$. Assume that this group is Z_p, and that G has charge q. Then there is some lowest power G^a ($a \geq 2$) which is invariant, and there can be terms in the Kahler potential of the form

$$\delta K = \frac{G^{a+1}\bar{G} + h.c.}{M^a}, \tag{4.1}$$

where M might be the unification scale or the Planck scale. This will give a high energy contribution to the would be axion mass[7]

$$(m_a^{high})^2 \sim \frac{M_P \Lambda^{1/4}}{M_1} (\frac{M_1}{M})^{a/2}. \tag{4.2}$$

The QCD contribution is

$$m_a^{QCD} \sim \frac{(100\mathrm{MeV})^2}{M_1} \tag{4.3}$$

For $M_1 \sim 1$ TeV the ratio of the high energy contribution to the QCD contribution is $\sim 10^9 (\frac{M_1}{M})^{a/2}$. Even for $a = 2$ and $M \sim 10^{15}$ GeV, the QCD contribution dominates. So we appear to have an axion but in an experimentally forbidden range.

There are three ways out of this problem. The simplest is to assume that only $\mathcal{F}$ and not the full U_A group is a classical symmetry of the $\mathcal{G}$ Lagrangian. That is, other terms in the superpotential for chiral fields charged under $\mathcal{G}$ break U_A down to a discrete subgroup containing $\mathcal{F}$. Alternatively, the classical U_A symmetry could be broken by a $\mathcal{G}$ anomaly. In either case, strong dynamics at the scale M_1 would give both scalar components of G masses of order $\frac{M_P \Lambda^{1/4}}{M_1}$.

A more interesting possibility is that U_A is a symmetry, but that our visible QCD axion actually evades the conventional experimental bounds. At the effective Lagrangian level it appears that the axion to goldstino pair amplitude is larger by a factor of $(\frac{\alpha_i}{\pi})^{-1}$ (where α_i is a standard model gauge coupling) than any decay into visible products. If the visible branching ratios are small enough, the conventional bounds might be evaded. A note of caution here is that the matrix element of the

[7]There can be other contributions coming from the superpotential, but these give smaller contributions to m_a.

leading operator mediating the decays into goldstinos may be chirally suppressed on shell[8]. Non-leading operators would contribute to p-wave decays into goldstinos and would have a suppression factor of order $(m_a/M_1)^2$ relative to the estimate above. This more than makes up for the gauge coupling suppression of visible decays. If this is indeed the case, the visible decays would dominate and such an axion is ruled out. We would then have to invoke breaking of U_A at the scale M_1 to construct a viable model.

In fact, our model *may* contain an alternative solution of the strong CP problem. To obtain it, we must make *another* assumption about the elusive $\mathcal{G}$ gauge theory: it should have automatic CP conservation. That is, the exact discrete R and $\mathcal{F}$ symmetries, the $\mathcal{G}$ gauge symmetry and renormalizability should guarantee the existence of an accidental CP symmetry of the $\mathcal{G}$ Lagrangian, under which G goes into its complex conjugate. In particular, in order to shift away the topological term in the $\mathcal{G}$ gauge Lagrangian, we rotate by the $U(1)$ R transformation under which all gauginos rotate, and all chiral superfields have R-charge 0. If the chiral multiplets in the theory fall into K irreducible $\mathcal{G}$ multiplets, the gauge interactions in the model are invariant under $K-1$ $\mathcal{G}$-anomaly free $U(1)$ symmetries, which are linear combinations of the phase rotations of the individual multiplets. If there are several multiplets in the same $\mathcal{G}$ representation, they are also invariant under $\mathcal{G}$-anomaly free $SU(m)$ transformations. The standard model gauge group is a subgroup of this anomaly free group, which also leaves all Yukawa couplings invariant[9]. The condition for CP invariance is that the full anomaly free group can transform away all phases in the Yukawa couplings that are allowed by $\mathcal{G} \otimes SU(3,2,1) \otimes R \otimes \mathcal{F}$.

One consequence of this assumption for the effective theory below M_1 is that the Kahler potential $G\bar{G}K(G/M_1, G^*/M_1)$ is CP invariant. We will also make the technically natural assumption that the minimum of the potential, which is derived from K once we add the superpotential to the limiting Lagrangian, is CP conserving; *i.e.* $< G >$ is real.

Furthermore, CP invariance of the strong $\mathcal{G}$ dynamics guarantees that the coefficients ϵ_i in the couplings $\epsilon_i (G/M_1)^a W_i^2$ to the standard model gauge bosons, are all real, and do not shift the value of θ_{QCD} when the VEV of G is turned on.

Now consider the low energy Lagrangian for the standard model coupled to G, still in the limiting model with $\Lambda = 0$. We can use the a combination of the $U(1)_R$ of gaugino rotation, the $U(1)_A$ which rotates all quark and anti-quark superfields by the same phase, and an equal phase rotation of the Higgs superfields[10]to eliminate

[8] I would like to thank S. Thomas for pointing this out to me.

[9] Recall our assumption that the model contains no relevant operators allowed by the symmetries.

[10] The independent linear combination of phase rotations on H_u and H_d is the gauged weak hypercharge. We will always choose a gauge in which the VEVs of the two Higgs fields have equal

both θ_{QCD} and argdet$g_u g_d$, (where $g_{u,d}$ are the up and down quark Yukawa coupling matrices), and to make the phase of the coupling g_μ in $g_\mu \int d^2\theta G H_u H_d$ real and negative.

In a nutshell, what we have shown is that, like the standard model before the discovery of the $U(1)_A$ anomaly, the limiting $\Lambda = 0$ model has all CP violation concentrated in the usual Jarlskog parameter of the CKM matrix.

Now consider what happens when $\Lambda \neq 0$. The superpotential for G comes from Planck scale physics near the horizon. It has no apparent reason to be CP invariant. In particular, the coefficent $w'(0)$ which determines F_G might be complex. Write $w'(0) = |w'(0)|e^{ia}$. In the low energy effective Lagrangian, below the scale M_1, F_G appears linearly and quadratically. The quadratic terms have the form $F_G F_G^*$. Thus, the phase a appears only in the gaugino masses

$$m^i_{1/2}\tilde{g}^i\tilde{g}^i = r^i \frac{|F_G|}{M_1} e^{ia}\tilde{g}^i\tilde{g}^i \tag{4.4}$$

(r^i are real numbers of order $\frac{\alpha_i}{\pi}$) and in the "b - term"

$$m^2_{ud} h_u h_d = g_\mu |F_G| e^{ia} h_u h_d + c.c. \tag{4.5}$$

The latter term is the only term in the effective potential that depends on the overall phase $e^{i(a_u+a_d)}$ of the Higgs fields. It is minimized by $e^{i(a_u+a_d)} = e^{-ia}$.

When the Higgs VEVs are substituted in the Yukawa couplings, this generates a phase for the determinant of the quark mass matrix

$$\text{arg det}\mathcal{M} = -3ia \tag{4.6}$$

We can eliminate both the phase of the gluino mass and that of the quark determinant by doing a $U(1)_R$ rotation with angle satisfying $e^{2i\theta_R} = e^{-ia}$, and a $U(1)_A$ rotation with $e^{12i\theta_A} = e^{3ia}$. A particular solution of these equations is

$$\theta_R = -2\theta_A = -\frac{a}{2} \tag{4.7}$$

Now recall that the Dynkin index, which determines the $SU(3)$ anomaly of the $U(1)$ rotation of a single Weyl fermion is 3 for the adjoint of $SU(3)$ and $\frac{1}{2}$ for the fundamental. $U(1)_A$ combines the $U(1)$ rotations of 12 Weyl fermions so it shifts θ_{QCD} by twice as much as a rotation by the same angle in $U(1)_R$. *But we have found that the $U(1)_A$ rotation we need to eliminate the phase of the quark determinant is half as large and has the opposite sign of the rotation we need to eliminate the gluino mass !* Thus, there is is no net θ_{QCD}.

phase.

There are several issues which must be checked before concluding that this is a real solution of the strong CP problem. We have only examined the effect of the superpotential $w(G)$ on the low energy effective theory. In fact, this interaction exists in the effective theory at any scale below the Planck scale. It is not clear whether renormalization effects coming from this term in the Lagrangian above the scale M_1 can invalidate our argument.

We have also used $U(1)_R$ transformations, without regard to their effect on irrelevant perturbations of the low energy theory. In our first use of a $U(1)_R$, to analyze the strong $\mathcal{G}$ dynamics, these corrections would be proportional to powers of $\frac{M_1}{M_U} \leq 10^{-11}$ and do not effect the argument. However we used a second $U(1)_R$ transformation below the scale M_1. Here the irrelevant operators are scaled by $1/M_1$, probably multiplied by powers of α_i/π from the standard model gauge interactions above M_1. These could provide new sources of CP violation and one must check that they do not induce a neutron electric dipole moment which contradicts experiment.

We have given an argument to the effect that the value of θ_{QCD} at the weak scale vanishes in our model. One must also check that the renormalization of this parameter between the weak scale and the scale at which the neutron EDM is measured, is small. A general argument to this effect, valid for a large class of theories, was given in [8]. One must check that that argument applies to the present model.

Finally, there is the issue of whether the many constraints on the $\mathcal{G}$ theory are such that they force us to have very light or massless G-hadrons, which contradict experimental bounds. In the next section we will note that such hadrons could even spoil our mechanism for SUSY breaking, which would mean that our model could not actually be a low energy effective Lagrangian for CSB. As far as I can see, the constraint of solving the strong CP problem does not add to this worry, but in the absence of a specific model, it is hard to be certain.

5. In search of a microscopic model

With the exception of the remarks about dark matter in the penultimate subsection, the phenomenological properties of our model can be expressed in terms of a lagrangian involving only the fields of the SUSic Standard Model and the Goldstino field, G. That Lagrangian is however non-renormalizable, and several of its crucial properties (*e.g.* the sign of the quartic term in the Kahler potential) depend on physics at the scale M_1 and above. Furthermore, at least one of the scalar components of the G field gets a mass close to M_1 once R-symmetry breaking is taken into account. One feels a moral compulsion to present a UV completion of the model which is valid up to the unification

scale, and explains the details of the low energy effective lagrangian. So far, I have not come up with such a model.

The obvious candidate for the $\mathcal{G}$ theory is SUSY QCD with $N_F = N_C + 1$. This model is asymptotically free, and has a vacuum state preserving both SUSY and a chiral R symmetry. We can couple G to combinations of the gauge invariant operators $\bar{F}_b F^a$. If $N_F = 5$ it has an anomaly free $SU(5)$ flavor symmetry, into which we can embed the standard model. The Yukawa couplings of G should of course preserve the standard model gauge symmetry.

Unfortunately, this model does not break SUSY[11], even when supplemented by the superpotential $w(G)$. The superpotential depends on the composite degrees of freedom $M^a_b = \bar{F}_b F^a$ and the baryons $B_a = \epsilon_{a,a_1 \ldots a_4} F^{a_1}_{A_1} \ldots F^{a_4}_{A_4} \epsilon^{A_1 \ldots A_4}$, and $\bar{B}^a = \epsilon^{a,a_1 \ldots a_4} \bar{F}^{A_1}_{a_1} \ldots \bar{F}^{A_4}_{a_4} \epsilon_{A_1 \ldots A_4}$. For small G it has the form

$$W = F_G G + M_1 G Y^b_a M^a_b + y_B B_a \bar{B}^b M^a_b + y_D \mathrm{det} M, \tag{5.1}$$

where F_G was defined above. This superpotential has SUSic minima, a single point with $B = \bar{B} = 0$, and a baryonic branch.

The model possesses several features whose generality is to be feared. Most asymptotically free gauge theories which can couple to the standard model will have continuous chiral symmetries with standard model anomalies. Indeed, anomaly matching is likely to be a key argument in showing that the model does not break R symmetry. This means that the model predicts additional massless degrees of freedom. In particular, the operator to which G couples may often be a free massless field M at low energies, in the theory without the coupling to G. The G coupling then provides a mass term and the expectation value of the free field M can cancel the F term that comes from interactions with the horizon. It is thus a significant challenge to produce a microscopic model which accomplishes our goals. On the other hand, if these arguments do not lead to a no-go theorem, we can hope that the low energy dynamics will be highly constrained.

6. Conclusions

I have presented an effective field theory of SUSY breaking which is based on the the idea of CSB. At the effective level, it contains one new singlet chiral superfield G. When the c.c. is set to zero, the theory is exactly super-Poincare invariant. It is also invariant under a discrete complex R symmetry as well as an ordinary discrete symmetry $\mathcal{F}$. The R charge of G is zero, and if there are no low energy fields with R

[11] I would like to thank N. Seiberg for helping me to analyze this system.

charge 2, this guarantees that G is massless. R symmetry charges of standard model fields are chosen to allow all of the SSM couplings while forbidding all dimension 4 and 5 operators that violate baryon and lepton number, except for dimension 5 operators that generate neutrino masses.

G has a renormalizable $g_H \int d^2\theta G H_u H_d$ coupling to the unique gauge invariant dimension 2 operator in the SSM. $\mathcal{F}$ symmetry charges are chosen to ensure that the μ term of the SSM can be generated only by the VEV of G[12].

SUSY is broken, and the VEV of G determined, by a superpotential of the form $M_P^2 \Lambda^{1/4} w(G/M_P)$, combined with a non-renormalizable Kahler potential $K = GG^* k(G/M_1, G^*/M_1)$, with $M_1 \ll M_P$. The form of the superpotential follows from the hypothesis of CSB. The constraint on the magnitude of M_1 is chosen for blatantly phenomenological reasons. There is then a SUSY breaking minimum with $|G| \sim M_1$. The constant term in the expansion of w around $G = 0$ can be used to tune the effective cosmological const to its observed value. The VEV of G gives rise to a μ term for the SSM, while the VEV of $F_G \sim M_P \Lambda^{1/4}$ gives rise to the μ_{ud}^2 scalar mass $h_u h_d + h.c.$.

Gaugino masses are generated by couplings $\epsilon_i \frac{\alpha_i}{\pi} (G/M_1)^a W_\alpha^2$[13]. In order that the various low energy SUSY breaking parameters be within experimental bounds, we must choose $M_1 \sim 1$ TeV. This means that we are straining the bounds of effective field theory, and it behooves us to construct a more microscopic model of physics at the scale M_1. This has proved to be difficult. Let me summarize the constraints on the microscopic theory.

- It must generate the dynamical scale M_1.
- It must not break SUSY, R or $\mathcal{F}$, explicitly or spontaneously.
- It must have a marginal coupling $g_{\mathcal{G}} \int d^2\theta G O_2$, where O_2 is a dimension two operator of R charge two, which is consistent with all of the symmetries of the problem. The non-trivial Kahler potential for G, which fixes $< |G| > \sim M_1$, is generated via this coupling. Note that we need the $\mathcal{F}$ symmetry to explain either the absence of the SSM μ term, or a term involving O_2 without G (either one of these could be shifted into the VEV of G). It is important that the operator O_2 not appear as a dimension one field in the effective theory below M_1. This is the property that we have found hard to realize in explicit models.

[12]At this level of analysis, we could simply absorb the μ term in G, but when terms in the $\mathcal{G}$ lagrangian are taken into account, we must impose a symmetry to eliminate it.

[13]We can also have couplings of the form $\ln(G/M_1) W_\alpha^2$ if U_A symmetry is unbroken by strong dynamics at the scale M_1. This leads to a QCD axion, which is probably ruled out by experiment.

- It might have an accidental $U(1)$ symmetry, U_A, with standard model anomalies, which can explain a logarithmic form of the non-renormalizable couplings of G to the standard model gauge fields. $\mathcal{F}$ is an anomaly free discrete subgroup of U_A. If this accidental symmetry is unbroken by either classical or quantum effects at the scale M_1 the model has a QCD axion. The axion decay constant is in a range that is probably ruled out by experiment, unless the axion decay mode into gravitinos dominates its visible decays.

- We have found a tentative solution to the strong CP problem without an axion, if the strong $\mathcal{G}$ dynamics is automatically CP conserving.

- We also need an accidental symmetry of the theory at scale M_1, which will prevent the generation of the dimension five operator $\int d^2\theta H_u^2 LL$ with coefficient $\frac{1}{M_1}$, but permit it with a coefficient an order of magnitude below the unification scale. This accidental lepton number, could conceivably be a subgroup of U_A larger than $\mathcal{F}$.

The model also contains tantalizing hints of connections to other important problems in particle physics. We have mentioned the remote possibility of a viable QCD axion. We have also noted that the discrete symmetry $\mathcal{F}$ could play the role of a horizontal symmetry, if we allow it to be generation dependent. In order to implement this, one must assume that $|G|/M_1$ is small, perhaps of order .2, the Cabibbo angle. Note that, given the formula for gaugino masses, this could also raise the scale M_1 and make our effective field theory approximation a little more palatable, but only if the G coupling to standard model gauge fields has the WZ form. This implies the low scale axion, which is probably ruled out. If we instead have couplings of the form $(G/M_1)^a W^2$, a small value of $< G >$ would force the scale M_1 to be very low and the model is ruled out. In addition to this problem, conversations with Y. Nir have convinced me that implementation of the idea that $\mathcal{F}$ is a horizontal symmetry is likely to remove the attractive flavor properties of the model. Indeed, as it stands, the only terms in the lagrangian which break the large flavor symmetry of the standard model gauge theory, are the Yukawa couplings to the Higgs boson. There is a consistent picture in which all issues having to do with flavor and neutrino masses, are fixed at scales within an order of magnitude below the GUT scale. The model then contains a natural GIM mechanism and there is no SUSY flavor problem.

Both of these issues should be explored more thoroughly. Additional work is also necessary to determine whether the mechanism for $SU(2) \times U(1)$ breaking in this model requires tuning. There are extra contributions to the quadratic term in the Higgs potential (beyond those familiar from gauge mediation), which involve the strong

coupling dynamics of the theory at scale M_1. We may have to wait for an explicit model of this sector before we can assess the answer to this question.

Some readers may be disturbed by the near coincidence between the scale M_1 and the logically independent scale $\sqrt{M_P \Lambda^{1/4}}$. We must postulate this coincidence for phenomenological reasons, but there is no apparent dynamical reason for it in the low energy model. I would love to find such a dynamical mechanism, but I am not sure that the coincidence is so much worse than that between the weak scale and the QCD scale (which some authors have found puzzling enough to require an explanation). According to the tenets of CSB, R symmetric couplings, like the $\mathcal{G}$ gauge coupling at the unification scale, are determined (up to very small corrections) by their values in the limiting SUSic theory. If the limiting SUSic model is unique, and the value of this coupling (which determines M_1) happened to work out correctly, we would consider it a great triumph. So, the puzzle of the coincidence between M_1 and the SUSY breaking scale may be, like flavor and neutrino masses, a puzzle that will only be resolved when we learn the full high energy theory.

To summarize this summary, I have found what appears to be an attractive model of SUSY breaking, motivated by the ideas of CSB. It solves many of the problems of other approaches, and presents us with a new candidate solution of the strong CP problem. More work is necessary to completely assess its phenomenological viability, and to work out its detailed predictions for physics at the TeV scale.

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